

Signatures of Baryon non-conserving Yukawa couplings in a supersymmetric theory

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Abstract

Renormalization effects of large baryon-nonconserving Yukawa couplings $\lambda''_{ijk} \overline{U_i D_j D_k}$ lower the right handed squark masses keeping the left-handed squark masses virtually untouched at the lowest order. At low energy they enhance the mass-splitting between left and right handed squarks of the same generation as well as intergenerational mass splitting among squarks, potentially detectable in future colliders or in rare decays. The predicted mass of the lightest stop squark becomes lower than the experimental bound for larger ranges of parameter space than that of the Baryon-conserving case, hence, further constraining the parameter space of a supersymmetric theory when baryon violation is included.

We know that the radiative corrections due to a large top quark Yukawa coupling drives a squared Higgs scalar mass to negative values triggering the radiative electroweak symmetry breaking [1] in the Minimal Supersymmetric Standard Model (MSSM). Can similar effects be caused by other large Yukawa couplings? In particular, MSSM allows the presence of lepton and baryon number violating couplings [2] unless they are forbidden by fiat invoking R-parity. In the models with explicit R-parity violation some of these couplings, especially the ones involving the third generation fields, are almost unconstrained by experiments. They could as well be as large as the top quark Yukawa coupling. Such Yukawa couplings may involve (for example) colored fields. Hence, intuitively we may expect that the radiative corrections due to such potentially large couplings can have analogous effects on squark masses, thereby serving as an indirect signature of R-parity violation. Here we explore this possibility.

Apart from the standard Yukawa couplings related to the fermion masses, Gauge invariance and Lorentz invariance allows the R-parity violating Yukawa couplings in the low energy effective supersymmetric theory, given by,

$$W_R = \frac{\lambda_{ijk}}{2} L_i L_j E_k^c + \lambda'_{ijk} L_i Q_j D_k^c + \frac{\lambda''_{ijk}}{2} D_i^c D_j^c U_k^c. \quad (1)$$

The simultaneous presence of the lepton and baryon number violating couplings give rise to fast proton decay unobserved in nature [3]. This leaves us the choice of allowing either the baryon number violating couplings (λ'') or the lepton number violating couplings (λ, λ') [4]. There are existing experimental bounds on most of the lepton number violating couplings [5], and so, we will consider the first scenario allowing only the presence of λ'' couplings which stand relatively unconstrained from experimental results as well as cosmological bounds.

The color antisymmetry in the first and second indices of λ''_{ijk} leaves us nine independent couplings. All of them were believed to have strong cosmological constraints [6] due to the conjecture that they would erase the primordial baryon asymmetry [7]. However, it has been shown [8] that such constraints are model dependent and λ'' can be completely free of cosmological bounds in scenarios with GUT-era leptogenesis. The laboratory bounds on λ''_{1j1} are derived from the non-observation of N - \bar{N} oscillation and double nucleon decay processes [9, 11]. The λ'' couplings involving the third generation twice, i.e., λ''_{133} and λ''_{233} are bounded from above by perturbative unitarity arguments [10] and they can very well be of order $O(1)$ at the unification scale. Moreover, we know that the third generation Yukawa couplings in the R-conserving sector are the largest ones; keeping this in mind we will have a similar hypothesis in the R-violating sector. We will consider the effects of λ''_{133} and λ''_{233} , where the third generation occurs twice, discarding the rest, on the renormalization of soft mass terms here.

The theoretical upper bounds [10, 11, 12] on λ''_{133} and λ''_{233} can be found demanding the perturbative unitarity of *all* the Yukawa couplings of the theory up to the unification scale M_X , taken to be 2×10^{16} GeV here. The essential point to get these bounds is that when λ''_{133} or λ''_{233} are large enough, they drive the top quark Yukawa coupling h_t into the non-perturbative domain before M_X , and they may also blow up along with h_t for low top quark masses, e.g., $m_t^{pole} = 165$ GeV in the low $\tan \beta$ region. We note that for a fixed m_t , the value of h_t depends on the ratio of the vevs of the two Higgs scalars $\tan \beta = \langle H_2 \rangle / \langle H_1 \rangle$. Hence, the largest allowed value for λ''_{133} and λ''_{233} depends on $\tan \beta$. We have taken¹ $\alpha_s(m_Z) = 0.114$,

¹The value of $\alpha_s(m_Z)$ is chosen to have unification in the gauge sector at the one-loop level.

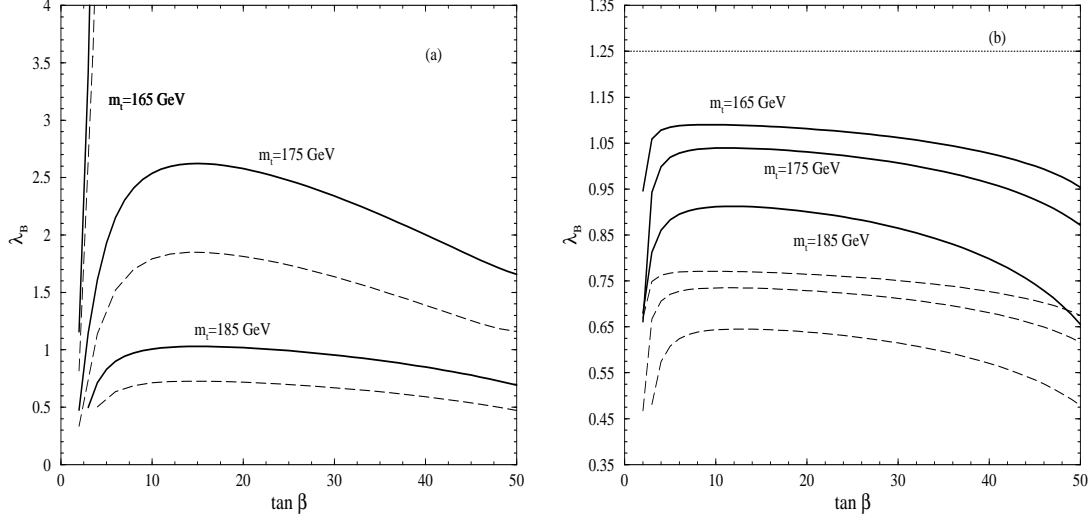


Figure 1: Largest allowed values of λ''_B (a) at the scale M_X and (b) at the scale m_t , for $m_b = 4.4$ GeV, $m_\tau = 1.777$ GeV, $\alpha_s(m_Z) = 0.114$. The dashed lines represent the case where $\lambda''_{133} = \lambda''_{233}$. The dotted horizontal line is the case where h_t , h_b and h_τ are neglected as in Reference [10].

$m_b(m_b) = 4.4$ GeV, $m_\tau(m_\tau) = 1.777$ GeV and several values m_t^{pole} . Using these values and the perturbative unitarity argument for *all* Yukawa couplings, we get the upper bounds on,

$$\lambda''_B = \sqrt{\lambda''_{133}{}^2 + \lambda''_{233}{}^2}. \quad (2)$$

These bounds are plotted in Figure (1) as a function of $\tan \beta$. We have not assumed $b - \tau$ unification to be as model independent as possible. Figure (1.a) shows the bounds at M_X , and Figure (1.b) those extrapolated to the scale m_t . To see the maximal renormalization effect on the squark masses for a given m_t^{pole} we have fixed $\tan \beta$ where λ''_B is at its maximum. For example, in our case²: $m_t^{pole} = 175$ GeV, $\tan \beta = 14.99$ and $\lambda''_B(M_X) = 2.62$ ($\lambda''_B(m_t) = 1.03$).

We have derived the evolution equations of the soft masses, including λ''_{133} and λ''_{233} [15]. In the case $\lambda''_B = 0$ the equations can be found in Ref. [1]. However, the essential differences

in the evolution equations can be traced from Eqns. (11), (12), (13) given below. We

²This perturbative bound is on λ''_B . Individual experimental bounds coming from LEP are $(\lambda''_{ij3}(m_Z) \leq 0.97)$ [13]. For the bounds on the product $\lambda''_{133}\lambda''_{233}$ see Reference [14].

have run the coupled set of renormalization group equations (RGE) for the soft masses and dimensionless Yukawa couplings numerically, taking into account the constraints coming from the radiative electroweak symmetry breaking. The unknown universal soft supersymmetry breaking parameters at the unification scale are,

$$m_0, m_{1/2}, A_0, A_B^0, \mu_0, B_0,$$

where we have taken $A_{133}^0 = A_{233}^0 = A_B^0 = A_0$ for simplicity, and we have chosen the sign of μ_0 to be positive. We have varied $m_{1/2}$ in the range 50 to 500 GeV, m_0 in the range 100 to 300 GeV, for the cases $A_0 = 0$ GeV and $A_0 = 100$ GeV. The other two parameters, μ_0 and B_0 , will be fixed by the choice of $\tan \beta = 14.99$ and the condition of electroweak symmetry breaking, given by,

$$B(\overline{Q}) = \left[\frac{(m_{H_1}^2 + m_{H_2}^2 + 2\mu^2) \sin 2\beta}{2\mu} \right]_{\overline{Q}}, \quad (3)$$

$$m_Z^2(\overline{Q}) = 2 \left[\frac{m_{H_1}^2 - m_{H_2}^2 \tan^2 \beta}{\tan^2 \beta - 1} - \mu^2 \right]_{\overline{Q}}, \quad (4)$$

where \overline{Q} is the scale [16] in which the minimum of the one-loop effective Higgs potential [17] becomes identical to the tree level one, given by,

$$\overline{Q} = \frac{1}{\sqrt{e}} \prod_i (m_i)^{d_i m_i^2 / \sum_p d_p m_p^2}, \quad (5)$$

and,

$$d_p = (-)^{2s_p} (2s_p + 1) \frac{\partial m_p^2}{\partial v_2}. \quad (6)$$

where m_p are *all* the mass eigenvalues at the scale \overline{Q} . When λ_B'' is turned on, $h_t(m_t)$ tends to decrease so $h_t(M_X)$ has to increase; however, the renormalization effects of the B-violating

couplings reduce the right handed squark masses (which is multiplied by h_t in the β -function), therefore $m_{H_2}^2$ increases (i.e., it becomes less negative), while $m_{H_1}^2$ remains almost the same. From Eqn.(4) we see that for a fixed value of $\tan \beta$, the value of $\mu(\overline{Q})$ has to decrease. This is achieved due to the increased top quark Yukawa coupling involved in the RGE for μ which gives the correct $h_t(m_t)$ to fit m_t^{pole} too. However, in order to avoid too much reduction in $\mu(\overline{Q})$, the initial value μ_0 has to increase.

Absence of the renormalization effect due to the R-parity violating couplings in the left handed sector creates an additional splitting between the squark masses. In order to get the physical masses of the squarks after electroweak symmetry breaking, one has to diagonalize the corresponding mass matrices. For example, in the case of the third generation up-type squarks the left-right mass matrix is [17],

$$\mathcal{M}_u = \begin{pmatrix} M_L^2 & m_u [A_u + \mu \cot \beta] \\ m_u [A_u + \mu \cot \beta] & M_R^2 \end{pmatrix}, \quad (7)$$

where,

$$M_{L/R}^2 = \tilde{m}_{L/R}^2 + m_u^2 + (\text{D - terms})_{L/R}. \quad (8)$$

For the down-type squarks $\cot \beta$ is replaced by $\tan \beta$. For the first and second generation down type squarks, off-diagonal elements are generated by renormalization effects of λ_{133}'' and λ_{233}'' , giving the 2×2 right-right mass matrix,

$$\mathcal{M}_{12}^{RR} = \begin{pmatrix} M_{dd}^2 & \tilde{m}_{ds}^2 \\ \tilde{m}_{ds}^2 & M_{ss}^2 \end{pmatrix}, \quad (9)$$

where,

$$\mathcal{M}_{ii} = \tilde{m}_{iR}^2 + (\text{D - terms}). \quad (10)$$

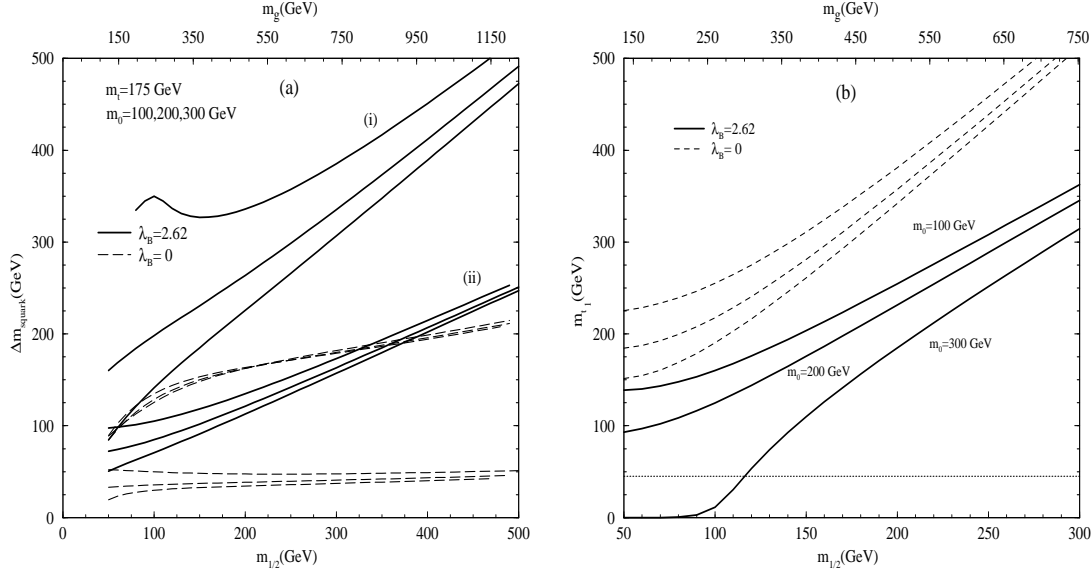


Figure 2: (a) The splitting in the (i) stop, (ii) sbottom, for several values of m_0 . In each group, the lowest curve is for the lowest value of m_0 and so on. Bold solid lines represent $\lambda_B''(M_X) = 2.62$, and dashed lines represent $\lambda_B'' = 0$, and in both cases $\tan\beta = 14.99$. (b) The lower eigenvalue of the stop system. Same convention of solid and dashed lines as before. The horizontal line is at 45 GeV, the lower bound from LEP. For $m_0 = 300$ GeV the lower stop may become lighter than the LEP bound for $m_{1/2} < 120$ GeV. We note that in $\lambda_B'' = 2.62$ case the lowest stop mass reduces with increasing m_0 , contrary to the $\lambda_B'' = 0$ case.

All the entries in the mass matrices are running parameters. We have shown the difference in the eigenvalues of the stop and sbottom squarks evaluated at their own scale in Figure (2.a). The dashed lines are for the case $\lambda_B'' = 0$. The splitting between the squark eigenvalues increases with the increasing $m_{1/2}$ as well as with increasing m_0 . The mass of the lighter stop is plotted in Figure (2.b). We notice that for large values of m_0 and low values of $m_{1/2}$ it can be below the experimental bound. In Figure (3) we have plotted the difference in mass eigenvalues of the first and the second generation right handed squarks. Changing the value of A_0 from 0 to 100 GeV we have not found any substantial difference in these results with B-violation shown in the solid lines. However, such a change in A_0 has observable effects in the difference in the stop eigenvalues when $\lambda_B'' = 0$.

The additional splitting created between the squark eigenvalues in the $\lambda_B'' \neq 0$ case is

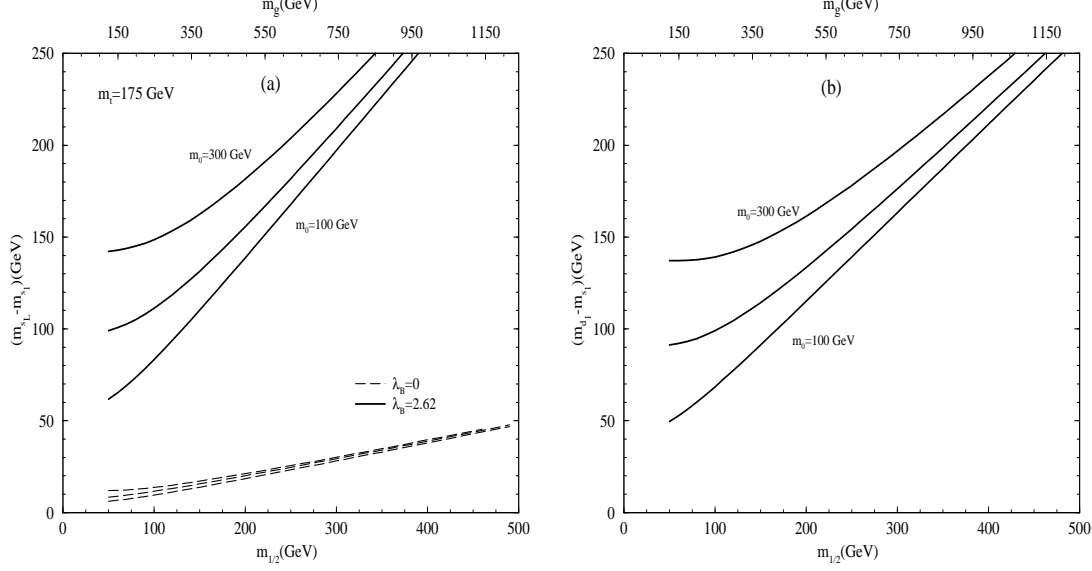


Figure 3: (a) Splitting between the left-handed s-strange squark and the lighter eigenvalue of the mass matrix in Eqn. (9), for several values of m_0 ; dashed lines are as before. (b) The difference in eigenvalues of the first and second generation right handed squark mass matrix. For the analytical expression see Eqn. (14) and Eqn. (15). Subscript 1 refers to the lighter eigenvalue as in the stop system.

mainly due to the induced splitting between the running masses in the diagonal entries, and it can be understood if we write down the RGEs for the differences in the left and right-handed running mass parameters. We have defined $Y_k = h_k^2/4\pi$, where h_k is a relevant Yukawa coupling and $t = \ln \mu/2\pi$, where μ is the renormalization scale and we have used the notation $Y_i = Y_{i33}$. The differences in the diagonal entries renormalize as,

$$\frac{\partial(\tilde{m}_{t_L}^2 - \tilde{m}_{t_R}^2)}{\partial t} = -\left(3\alpha_2 M_2^2 + Y_t \Delta_u + 2Y_1 \Delta_1 + 2Y_2 \Delta_2 + 4\sqrt{Y_1 Y_2} \tilde{m}_{ds}^2\right) + \alpha_1 M_1^2 + Y_b \Delta_d, \quad (11)$$

$$\frac{\partial(\tilde{m}_{b_L}^2 - \tilde{m}_{b_R}^2)}{\partial t} = -\left(3\alpha_2 M_2^2 + Y_b \Delta_d + 2Y_1 \Delta_1 + 2Y_2 \Delta_2 + 4\sqrt{Y_1 Y_2} \tilde{m}_{ds}^2\right) + \frac{1}{5}\alpha_1 M_1^2 + Y_t \Delta_u, \quad (12)$$

$$\frac{\partial(\tilde{m}_{s_L}^2 - \tilde{m}_{s_R}^2)}{\partial t} = -\left(3\alpha_2 M_2^2 + 2Y_2 \Delta_2 + 2\sqrt{Y_1 Y_2} \tilde{m}_{ds}^2\right) + \frac{1}{5}\alpha_1 M_1^2 \quad (13)$$

In the right handed mass matrix, the difference in eigenvalues can be computed from,

$$\frac{\partial(\tilde{m}_{d_R}^2 - \tilde{m}_{s_R}^2)}{\partial t} = 2(Y_1 \Delta_1 - Y_2 \Delta_2), \quad (14)$$

$$\frac{\partial(\tilde{m}_{ds}^2)}{\partial t} = (Y_1 + Y_2) \tilde{m}_{ds}^2 + \sqrt{Y_1 Y_2} (\Delta_1 + \Delta_2). \quad (15)$$

with,

$$\begin{aligned}
\Delta_u &= A_t^2 + \tilde{m}_{q_{L3}}^2 + m_{H_2}^2 + \tilde{m}_{u_{R3}}^2, \\
\Delta_d &= A_b^2 + \tilde{m}_{q_{L3}}^2 + m_{H_1}^2 + \tilde{m}_{d_{R3}}^2, \\
\Delta_i &= A_{i33}^2 + \tilde{m}_{u_{R3}}^2 + \tilde{m}_{d_{R3}}^2 + \tilde{m}_{d_{Ri}}^2,
\end{aligned} \tag{16}$$

When $\lambda''_{\vec{P}}$ is turned on in the R.H.S. of Eqn.(11), (12), (13) the splitting between the squarks is increased. This also happens with increasing $m_{1/2}$ due to the presence of M_2 in the right hand side. Large values of $m_{1/2}$ also increase the Δ 's and this gives a indirect additional effect to amplify the splittings. The Δ 's also grow with m_0 and hence the splitting is also enhanced by the rise in m_0 . The flattening of the curves in the $\lambda''_{\vec{P}} = 0$ case in Fig. (2.a) is due to the off-diagonal elements in the mass matrices. In the $\lambda''_{\vec{P}} \neq 0$ case, both the values of μ and A_u reduce at low energy compared to the B-conserving case making the off-diagonal effect milder. In Eqn.(11) Y_t has an increasing effect (negative sign) while Y_b a decreasing effect (positive sign). Whereas in Eqn.(12) Y_b and Y_t exchange their roles. Because Y_t is much greater than Y_b , splitting in the stop system is more pronounced.

The absence of Flavor Changing Neutral Current (FCNC) constrains the off-diagonal mixing terms between the first and the second generation [18] denoted by \tilde{m}_{ds}^2 in Eqn. (9). Even though at the GUT scale the mixing is taken to be vanishing, it is generated by one loop diagrams connecting the right handed d and s squarks when λ''_{133} and λ''_{233} are present simultaneously. However this diagram is suppressed when one coupling is sufficiently smaller than the other [15]. Now we will show, that the maximum mass-splittings plotted in Figures

(2) and (3) are independent of the relative magnitude of λ''_{133} and λ''_{233} . In other words, even after taking into account the constraints from FCNC on the ratio $\lambda''_{133}/\lambda''_{233}$, the mass splittings remain unaltered for given value of the sum $\sqrt{\lambda''_{133}{}^2 + \lambda''_{233}{}^2}$.

Let us re-write the splitting in the stop system in terms of the quantities,

$$r^2 = Y_1/Y_2 \quad \text{and} \quad c(r) = \Delta_1 Y_1 + \Delta_2 Y_2 + 2\sqrt{Y_1 Y_2} \tilde{m}_{ds}^2. \quad (17)$$

The Eqn. (11) is of the general form,

$$\frac{\partial(\tilde{m}_{t_L}^2 - \tilde{m}_{t_R}^2)}{\partial t} = f - 2 c(r), \quad (18)$$

where, f is independent of r and the combination c renormalize as,

$$\frac{\partial c(r)}{\partial t} = g + h c(r), \quad (19)$$

g and h being independent of r . Differentiating Eqn. (19) with respect to r and performing the integration of Eqn. (19) from the scale M_X to a scale μ we get,

$$\left. \frac{\partial c}{\partial r} \right|_{\mu} = \left. \frac{\partial c}{\partial r} \right|_{M_X} e^{\int_{M_X}^{\mu} h dt}. \quad (20)$$

Universality of the soft masses at the scale M_X gives $\left. \frac{\partial c}{\partial t} \right|_{M_X} = 0$. Hence $\partial c/\partial r$ vanishes at all scales proving c to be independent of r . As Eqn. (11) is independent of r , so is the splitting for a given value of λ''_B . For the first two generations the eigenvalues of the right-handed-squark mass matrix can be shown to be independent of r by repeated application of this procedure to the evolution of the trace and the determinant of the mass-matrix given in Eqn. (9).

To conclude, non zero $\lambda''_{\vec{B}}$ introduces asymmetric renormalization effects between the left and right handed sector of squarks, increasing the splitting between them, as shown in Figure (2) and (3), which has been obtained by full numerical integration of the all soft masses and the third generation dimensionless Yukawa couplings of MSSM extended by including baryon non-conserving Yukawa couplings of the third generation, and properly taking into account the constraints imposed by correct radiative symmetry breaking. We have shown that these splittings are independent of $\lambda''_{133}/\lambda''_{233}$ but depends on $\lambda''_{\vec{B}}$. The pattern that emerged here is that the splitting grows with increasing m_0 and $m_{1/2}$. Our solid lines for the stop system in Figure (2.a) can be interpreted as the maximum possible splitting (corresponding to $\lambda''_{\vec{B}}(M_X) = 2.62$ or $\lambda''_{\vec{B}}(m_t) = 1.03$), and similarly for the sbottom and s-strange systems.

The lower eigenvalue in the stop and the sbottom system is considerably lower than the $\lambda''_{\vec{B}} = 0$ case. The lightest stop mass is predicted to be below the experimental bound of 45 GeV for ranges of parameter space when $\lambda''_{\vec{B}}$ is large. Hence, when B-violating couplings are included, wider ranges of parameter space of the supersymmetric theory can be eliminated from the experimental lower bound on squark masses. Contrary to the naive expectations, the mass of the lightest stop *reduces* with increasing m_0 .

The searches of R-parity violating supersymmetry in colliders [19] have to take into account the susy mass spectrum with the proper renormalization effects due to R-parity violating couplings. On the other hand, the splitting in the stop system would also influence chargino contribution to the flavor changing decay $b \rightarrow s + \gamma$ [20].

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After the completion of this work, we came across the preprint by B. de Carlos and P. L. White, hep-ph/9602381, OUTP-96-01P, where RGE for the soft terms including R-parity violating Yukawa couplings have been derived. Our RGEs agree with them, however they have concentrated on leptonic flavor violating effects induced by L violating couplings.

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λ''_{133} and λ''_{233} removes the degeneracy between the squarks to a large extent [see Figure 2].

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